

A natural Z' model with inverse seesaw and leptonic dark matter

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Abstract

We consider a model for a Z' -boson coupled only to B-L and hypercharge. Besides the usual right-handed neutrinos, we add a pair of fermions with a fractional lepton charge, which we therefore call leptinos. One of the leptinos is taken to be odd under an additional Z_2 charge, the other even. This allows for a natural (inverse) seesaw mechanism for neutrino masses. The odd leptino is a candidate for dark matter, but has to be resonantly annihilated by the Z' -boson or the Higgs-boson responsible for giving mass to the former. Considering collider and cosmological bounds on the model, we find that the Z' -boson and/or the extra Higgs-boson can be seen at the LHC. With more pairs of leptinos leptogenesis is possible.

I. INTRODUCTION

The standard model (SM) gives an excellent description of the known laws of particle physics. However there are a few facts that it cannot explain. Of importance for this paper are the neutrino mass spectrum, in particular the smallness of neutrino masses, the baryon-antibaryon asymmetry in the universe and the presence of dark matter.

Though it appears to be very difficult to explain the mass-spectrum of the fermions, there is a mechanism that can in principle make it plausible why the neutrino masses are much smaller than the other masses. The mechanism uses the presence of right-handed neutrinos with a large Majorana mass. The mass-matrix also contains a Dirac-mass. If the Majorana mass is much larger than the Dirac mass, one finds after diagonalization one (very) light and one heavy neutrino. The method is called the seesaw mechanism. Several realizations exist, the simplest being the type-I seesaw [1–5]. In this paper we consider a somewhat more involved form, which is conventionally called the inverse seesaw mechanism in the literature [6, 7].

The presence of heavy neutrinos does not only affect the mass-spectrum, it can also provide a mechanism to explain the baryon asymmetry of the universe. The method is called “leptogenesis” [8] (also see [9–12] for recent reviews). An asymmetry in the lepton sector is produced via CP-violating heavy lepton decays. Subsequently the asymmetry in the lepton sector is transferred to the baryons by means of electroweak (EW) sphaleron processes [13–15]. Most early papers used neutrino masses at the grand unified scale, however a successful TeV scale leptogenesis is possible as well if two such neutrinos are degenerate in mass, thereby enhancing the CP -asymmetry parameter. This goes under the name of “resonant” leptogenesis [16–21]. In this case, flavour effects have to be taken into account [22–26].

Cosmological observations imply the existence of non-baryonic dark matter that drives structure formation on large scales and dominates galactic and extra-galactic dynamics. Within the SM there is no candidate for this matter. One therefore has to enlarge the SM. The easiest way to explain the dark matter is to postulate the existence of thermally produced Weakly Interacting Massive Particles (WIMPs), with masses roughly at the weak scale. An additional unbroken symmetry is postulated, that prevents the WIMPs from decaying.

The latest experimental searches (see, e.g., [27]) confirm the SM again at higher energy

scales than before. The data leave little space for modifications. In particular complicated extensions of the SM lead to phenomenological problems, for instance with flavour changing neutral currents, and tend to need many fine tunings of parameters. Minimal extensions are therefore preferable. The simplest form to enlarge the gauge group of the SM is to add a single $U(1)$ factor, which has to be a linear combination of hypercharge and B-L, baryon minus lepton number, if one does not want to enlarge the fermion spectrum. These are the so-called “minimal Z' models” (see, for example [28–33]). The spontaneous symmetry breaking of the extra $U(1)$ factor requires at least a new complex singlet scalar field. If the coupling of the new gauge group contains a term proportional to B-L, the absence of chiral anomalies demands the presence of additional SM-singlet fermions. The presence of right-handed neutrinos, one per generation, removes all anomalies [34–42].

In these minimal models a type-I seesaw mechanism can be introduced, dynamically generating neutrino masses. The parameters controlling the neutrino masses are compatible with resonant leptogenesis [43–46]. A Z_2 symmetry can be introduced to provide a stable DM candidate [47–50]. In these models with type-I seesaw, some fine-tuning is required to get two neutrinos almost degenerate and thereby have a successful TeV scale leptogenesis. The Z_2 symmetry is needed for stabilising the DM, but has no further relation with the neutrino masses.

The inverse seesaw appears more suitable to tackle these issues. Within this framework, 2 new neutrinos per generations are included. Once the mass matrix (of the left-handed and the 2 right-handed neutrinos) is diagonalized, besides the usual light SM-like neutrinos, 3 pairs of naturally quasi-degenerate heavy neutrinos appear, thereby easily implementing the requirements for resonant leptogenesis [51–53].

In the context of $U(1)_{B-L}$ extensions of the SM, a model with the inverse seesaw realization exists. The model contains two extra dileptons, one of which enters the neutrino mass matrix [54]. A Z_2 symmetry is advocated that leads to zeroes in the mass matrix.

A disadvantage of this model is that this Z_2 symmetry has to be broken, in order to avoid exactly zero-mass neutrinos. The breaking mechanism is not present in the tree-level Lagrangian, coming from non-renormalizable operators.

In this paper we discuss a consistent extension of the SM with a $U(1)$ gauge group, related to B-L number, providing for an inverse seesaw mechanism. In contrast to Ref. [54], we only use renormalizable operators. The mechanism is natural in the sense that we allow for

all renormalizable terms in the Lagrangian, consistent with the symmetries of the fields. We add pairs of fermions with fractional lepton number, so-called “leptinos” to the Lagrangian. One of them is odd under an additional Z_2 symmetry, the other even as are all ordinary SM particles. Both fermions are needed in order to cancel anomalies. The Z_2 symmetry, together with the B-L charge assignments, restricts the form of the neutrino mass matrix. At the same time it stabilises the odd leptino, that becomes the dark matter candidate in the model.

We will present in detail a version with only one extra pair of leptinos, which is sufficient for the discussion of dark matter. The possibility of successful leptogenesis requires the extension of the inverse seesaw mechanism to more generations, in order to provide the necessary large phases driving CP -violation. The detailed study is beyond the scope of this paper. However, we make some comments in the last section.

The paper is structured as follows. In the next section, the model is presented. Section III collects results for the dark matter abundance generated by the model. In section IV the possibility of a successful leptogenesis is outlined. Finally, in section V we present our conclusions.

II. THE MODEL

We base our extension of the SM on the minimal Z' model [29–33]. The SM gauge group is extended by including a $U(1)$ factor, related to the B-L number, with generic mixing with the $U(1)_Y$. A SM singlet complex scalar χ is required for the spontaneous symmetry breaking of the further $U(1)$ group, thereby providing the Z' boson a mass. The requirement of anomaly cancellation is fulfilled by introducing one right-handed (RH) neutrino per generation. Furthermore, the inverse seesaw mechanism needs extra SM singlet fermions, coming in pairs in order not to spoil the anomaly cancellation. Minimally, just one extra pair of fermions is sufficient to provide a DM candidate.

The classical gauge invariant Lagrangian, obeying the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ gauge symmetry, can be decomposed as:

$$\mathcal{L} = \mathcal{L}_{YM} + \mathcal{L}_s + \mathcal{L}_f + \mathcal{L}_Y. \quad (1)$$

A. Gauge sector

In the gauge field basis in which the kinetic terms in \mathcal{L}_{YM} are diagonal [33], the covariant derivative reads:

$$D_\mu \equiv \partial_\mu + ig_S T^\alpha G_\mu^\alpha + ig_W T^a W_\mu^a + ig_1 Y B_\mu + i(g_2 Y + g_{BL} Y_{B-L}) B'_\mu. \quad (2)$$

This generic model describes a continuous set of minimal $U(1)$ extensions of the SM, that can be labelled by the charge assignments of the particles [33]. As any other parameter in the Lagrangian, g_2 and g_{BL} are running parameters, therefore their values have to be set at some scale. Special sets of popular Z' models (see, e.g. [29, 55]) can be recovered by imposing relations between g_2 and g_{BL} . However such relations can be changed through the renormalization group running of the coupling constants.

It is instructive to consider the renormalization group equations (RGEs) for the gauge couplings. The one-loop RGEs read [28, 31, 56]

$$\frac{d}{dt} g_1 = \frac{1}{16\pi^2} [A^{YY} g_1^3], \quad (3)$$

$$\frac{d}{dt} g_{BL} = \frac{1}{16\pi^2} [A^{XX} g_{BL}^3 + 2A^{XY} g_{BL}^2 g_2 + A^{YY} g_{BL} g_2^2], \quad (4)$$

$$\frac{d}{dt} g_2 = \frac{1}{16\pi^2} [A^{YY} g_2 (g_2^2 + 2g_1^2) + 2A^{XY} g_{BL} (g_2^2 + g_1^2) + A^{XX} g_{BL}^2 g_2], \quad (5)$$

For the model we are discussing (Y is the SM weak hypercharge, $X = B-L$ is the B-L number), the coefficients are:

$$A^{YY} = \frac{41}{6}, \quad A^{XX} = \frac{32 + (Y_X^{B-L})^2}{3} + \frac{4}{27} N_\ell, \quad A^{YX} = \frac{16}{3}, \quad (6)$$

if N_ℓ generations of leptinos (section C) are included. Notice the small difference with respect to Ref. [31] in the A^{YY} coefficient, due to the SM Higgs in the counting. The equations can be solved algebraically [31]:

$$\frac{1}{g_1^2} + 2A^{YY} t = \text{constant}, \quad (7)$$

$$\frac{2A^{YY} g_2 + 2A^{YX} g_{BL}}{g_1^2 g_{BL}} = \text{constant}, \quad (8)$$

$$\frac{A^{YY} (g_1^2 + g_2^2) - A^{XX} g_{BL}^2}{g_1^2 g_{BL}} = \text{constant}, \quad (9)$$

Particularly interesting is eq. (8), which leads to an infrared (IR) fixed point for the model

$$41g_2 + 32g_{BL} = 0. \quad (10)$$

This fixed point is independent of the additional matter we consider in the model. The reason for this is that the extra fields we introduce (RH neutrinos, leptinos and the singlet scalar) are all singlets under the SM gauge group, hence not entering in the diagonal and mixed hypercharge coefficients A^{YY} and A^{XY} , that are the only terms appearing in eq. (8). In other words, this IR fixed point is a model independent property of the minimal Z' model.

For the following study, it is more important to focus on measurable observables, one of which being the Z' total width. In the approximation of massless fermions and neutrinos, with N_ℓ generations of leptinos, of which just the CP -odd leptinos (S_2 , the DM candidate) are massive, with in first approximation degenerate masses M_{S_2} , the total width reads

$$\Gamma_{Z'} = \frac{M_{Z'}}{12\pi} \left(\left(8 - \frac{N_\ell}{2} + \frac{1}{18} \sqrt{1 - \left(\frac{2M_{S_2}}{M_{Z'}} \right)^2} \right) g_{B-L}^2 + \left(5 - \frac{N_\ell}{8} \right) g_2^2 + \left(\frac{13}{2} + \frac{3 - N_\ell}{2} \right) g_2 g_{BL} \right). \quad (11)$$

Since the S_2 particle is only right-handed, its hypercharge is zero (see table II), so that its coupling to the Z' boson does not depend on g_2 . Moreover, in the evaluation of the DM relic abundance with S_2 as the DM candidate, we will see that a resonant annihilation with the Z' boson is required. In these conditions, $2M_{S_2} \sim M_{Z'}$ and $\text{BR}(Z' \rightarrow S_2 S_2) \rightarrow 0$. Being around the resonance, the only parameter that can influence the $S_2 S_2 \rightarrow Z'$ partial amplitude, and thus the relic density, is therefore the Z' boson width. The smaller the total width, the higher the relative partial amplitude. Hence, the minimum relic density, when one keeps $M_{Z'}$ fixed, can be obtained by minimizing the Z' width of eq. (11) with respect to g_2 . By direct computation, we obtain that the total Z' decay width is minimized for

$$g_2^{\min}(g_{BL}) = -2 \frac{16 - N_\ell}{40 - N_\ell} g_{BL}. \quad (12)$$

When $N_\ell = 1$ (the case discussed here), $g_2^{\min} = -10/13 g_{BL}$. Other important cases are for $N_\ell = 3$ (discussed in section IV in connection with leptogenesis), for which we obtain $g_2^{\min} = -26/37 g_{BL}$, and $N_\ell = 0$, where the minimum width is for the $\text{SO}(10)$ -inspired $U(1)_X$ model, $g_2^{\min} = -4/5 g_{BL}$.

So far, we have neglected that the Z' boson is in general mixed with the SM Z boson. Typical bounds from LEP-I measurements at the Z -boson peak require the mixing angle to be less than $\mathcal{O}(10^{-3})$ [57].

Further, from a combination of LEP-I and LEP-II data the ratio mass-over-coupling is bounded to be bigger than several TeV. In Ref. [32] the authors reanalyzed the LEP data

for a model with the same gauge sector as ours, while a specific bound for $g_2 = 0$ can also be found in Refs. [29, 58].

Before moving on to the results, we present here (see figure 1) the 95% C.L. exclusions at the LHC in the $g_{BL} - g_2$ plane, based on the CMS data at $\sqrt{s} = 7$ TeV for the combination of $4.7(4.9) \text{ fb}^{-1}$ in the electron(muon) channels [59]. ATLAS has as well published an analysis for $\sim 5 \text{ fb}^{-1}$ [60], but their limits are less tight than the CMS ones. Therefore, we will not present them here.

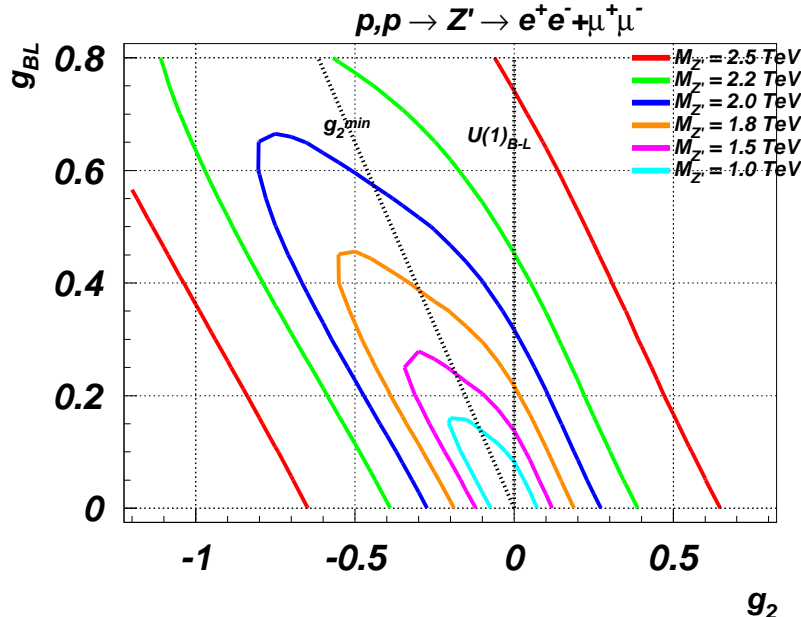


FIG. 1: Z' exclusions from CMS data, at $\sqrt{s} = 7$ TeV for the combination of 4.7 fb^{-1} in the electron channel and 4.9 fb^{-1} in the muon channel. The dotted black lines refer to the two main benchmark models of this analysis.

Table I collects the maximum allowed g_{BL} coupling per given Z' boson mass for the B-L model ($g_2 = 0$) and for the g_2^{min} case of eq. (12), for $N_\ell = 1$.

B. Scalar sector

In order to give the Z' boson a mass through spontaneous symmetry breaking, an extra complex scalar χ is introduced. The scalar Lagrangian reads

$$\mathcal{L}_s = (D^\mu H)^\dagger D_\mu H + (D^\mu \chi)^\dagger D_\mu \chi - V(H, \chi), \quad (13)$$

$M_{Z'}$ (TeV)	g_{BL} for g_2^{min}	g_{BL} for B-L
2.5	> 0.8	0.75
2.2	0.81	0.45
2.0	0.58	0.31
1.8	0.39	0.22
1.5	0.24	0.13
1.0	0.13	0.08
0.5	0.05	0.03

TABLE I: 95% C.L. exclusions for the benchmark models of interest. Couplings smaller than those in the table are allowed.

with the scalar potential given by

$$V(H, \chi) = m^2 H^\dagger H + \mu^2 |\chi|^2 + \lambda_1 (H^\dagger H)^2 + \lambda_2 |\chi|^4 + \lambda_3 H^\dagger H |\chi|^2, \quad (14)$$

where H and χ are the complex scalar Higgs doublet and singlet fields, respectively.

The scalar sector is now made of two real CP -even scalars, a light one h_1 and a heavy one h_2 , that are mixtures of the Higgs doublet field and the singlet field:

$$\begin{pmatrix} h_1 \\ h_2 \end{pmatrix} = \begin{pmatrix} \cos \alpha & -\sin \alpha \\ \sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} h \\ h' \end{pmatrix}, \quad (15)$$

where $v(v')$ is the VEV of the doublet(singlet) field, while the mixing angle $-\frac{\pi}{2} \leq \alpha \leq \frac{\pi}{2}$ can be expressed as:

$$\tan 2\alpha = \frac{\lambda_3 v v'}{\lambda_1 v^2 - \lambda_2 (v')^2}, \quad (16)$$

with $\lambda_{1,2,3}$ being the parameters entering in the quartic pieces of the scalar potential. The presence of extra heavy neutrinos and the Z' boson will alter the properties of the Higgs bosons. The scalar mixing angle α is a free parameter of the model, and the light(heavy) Higgs boson couples to the new matter content proportionally to $\sin \alpha(\cos \alpha)$, i.e. with the complementary angle with respect to the interactions with the SM particles.

The B-L-breaking vev v' can be expressed in terms of the Z' mass and the gauge couplings as follows:

$$v' = \frac{3M_{Z'}}{2g_{BL}} \sqrt{1 - \frac{g_2^2 v^2}{4M_{Z'}^2 - v^2(g^2 + g_1^2)}}. \quad (17)$$

Here we have chosen the B-L charge of the complex singlet scalar to be $Y_\chi^{B-L} = \frac{2}{3}$, as determined by our choice of the Yukawa sector.

The LEP bounds on $M_{Z'}/g_{BL}$ can provide an absolute lower bound on the VEV v' through eq. (17). For the benchmark scenarios of interest here, the bounds are:

$$g_2 = 0 : \quad v' > 9 \text{ TeV}, \quad (18)$$

$$g_2^{min} : \quad v' > 6.7 \text{ TeV}. \quad (19)$$

C. Yukawa sector

Moving to the Yukawa sector, the core of this paper, the inverse seesaw mechanism can be implemented by means of i extra pairs of RH fields beside the usual 3 RH neutrinos (ν_R), the latter required by the anomaly cancellation. With only renormalizable operators, the Yukawa Lagrangian for the neutrino sector reads

$$\mathcal{L}_Y^\nu = -y^D \overline{\ell_L} \tilde{H} \nu_R - y^N \overline{\nu_R^c} S_1^i \chi - y_i^{s1} \overline{(S_1^i)^c} S_1^j \chi^\dagger - y_i^{s2} \overline{(S_2^i)^c} S_2^j \chi + \text{h.c.}, \quad (20)$$

where S_2^i are the only odd fields under a Z_2 symmetry that is introduced to avoid unwanted S_1 - S_2 mixing. A summation over i, j is implied and an identical number of S_1 and S_2 fields is required by anomaly cancellation. As a concrete model for the DM study, we focus on $i = j = 1$ case. By convention, we choose to implement the inverse seesaw mechanism in the third generation of leptons only. In the following, the formulas will refer to the latter and we drop the “ i ” index, unless otherwise specified.

The gauge invariance of the Yukawa Lagrangian is achieved by solving $Y_{S_1}^{B-L} + Y_\chi^{B-L} = -Y_{\nu_R}^{B-L}$ and $2Y_{S_1}^{B-L} = \pm Y_\chi^{B-L}$, with Y^{B-L} the B-L charge of the fields. The following two sets of solutions exist, considering that $Y_{\nu_R}^{B-L} = 1$ is fixed by anomaly cancellations:

- i) $Y_{S_1}^{B-L} = -1$ and $Y_\chi^{B-L} = 2$, that is a replica of the type-I B-L model [38]. In this case, Majorana masses for all the neutrinos would be allowed, as well as a $\overline{\ell_L} \tilde{H} S_1$ term;
- ii) $Y_{S_1}^{B-L} = 1/3$ and $Y_\chi^{B-L} = 2/3$. This is the new solution proposed here, that forbids all unwanted terms. Because the B-L charge is one third of that of the normal leptons and because they do not carry a colour charge, the $S_{1,2}$ fields will be called “leptinos” throughout the rest of this paper.

ψ	ℓ_L	e_R	ν_R	S_1	S_2	H	χ
$SU(2)_L$	2	1	1	1	1	2	1
\mathbf{Y}	$-\frac{1}{2}$	-1	0	0	0	$\frac{1}{2}$	0
$\mathbf{B-L}$	-1	-1	-1	$\frac{1}{3}$	$-\frac{1}{3}$	0	$\frac{2}{3}$
Z_2	+	+	+	+	-	+	+

TABLE II: Y and $B-L$ quantum number and Z_2 parity assignments to chiral fermion and scalar fields.

Table II summarises the particle content and the charge assignments.

After spontaneous symmetry breaking, the two Higgs fields acquire vacuum expectation values (called v and v' , respectively). As a consequence, the first 3 terms of eq (20) lead to the following mass matrix for the interacting neutrinos, in the (ν_l^c, ν_R, S_1) basis:

$$\mathcal{M}^I = \begin{pmatrix} 0 & M_D & 0 \\ M_D^\dagger & 0 & M_N \\ 0 & M_N & M_{S_1} \end{pmatrix}, \quad (21)$$

where

$$M_D = \frac{y^\nu}{\sqrt{2}} v, \quad M_N = \frac{y^N}{\sqrt{2}} v', \quad M_{S_1} = \sqrt{2} y^{s1} v'. \quad (22)$$

When all 3 generations of leptinos are considered, y^ν and y^{s1} are in general 3×3 complex matrices, and it is not restrictive to consider y^N as a 3×3 real and diagonal matrix.

After diagonalization of eq. (21), the 3×3 neutrino mass matrices are

$$M_{\nu_l} \sim M_D M_N^{-1} M_{S_1} (M_N^T)^{-1} (M_D)^\dagger, \quad (23)$$

$$M_{\nu_h}^2 = M_{\nu_{h'}}^2 \sim M_D^2 + M_N^2, \quad (24)$$

However, it is beyond the scope of this paper to consider the most general case in detail. We limit ourselves to the model with only one generation of leptinos, which is largely sufficient for the dark matter question. In this case, one can without loss of generality choose a basis in which both the 3-component vector y^N and the single parameter y^{s1} are real. The above formulas simplify, since the seesaw mechanism acts only on the third generation, so that \mathcal{M}^I is a 3×3 matrix and eq. (23) gives the masses of the 3 neutrino eigenstates, where ν_h and $\nu_{h'}$ combine in a quasi-Dirac fermion. Regarding the first and second generations,

the LH and RH neutrinos obtain the usual Dirac mass term, for which $\mathcal{O}(10^{-12})$ Yukawa parameters are required to fit the light neutrino masses.

The mass matrix of eq. (21) is diagonalized by a 3×3 unitary matrix. Such a matrix can be parametrized by 3 angles and some phases. The latter will be neglected here. Formally, the following generic parametrization can be employed:

$$U = \begin{pmatrix} 1 & 0 & 0 \\ 0 & C_{23} & S_{23} \\ 0 & -S_{23} & C_{23} \end{pmatrix} \begin{pmatrix} C_{13} & 0 & S_{13} \\ 0 & 1 & 0 \\ -S_{13} & 0 & C_{13} \end{pmatrix} \begin{pmatrix} C_{12} & S_{12} & 0 \\ -S_{12} & C_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (25)$$

where $S_{ij}(C_{ij}) = \sin \alpha_{ij}(\cos \alpha_{ij})$. In good approximation, the angle defining the mixing between ν_R and S_1 is found to be very close to maximal, i.e., $\alpha_{23} \sim \pi/4$, while the angle parametrizing the mixing between ν_L and S_1 is very small: $S_{13} \sim M_D M_{S_1} / M_N^2 \ll 1$. The unitary matrix of eq. (25) can then be simplified to

$$U \simeq \begin{pmatrix} C_{12} & S_{12} & 0 \\ -\frac{S_{12}}{\sqrt{2}} + \frac{C_{12}}{\sqrt{2}} & \frac{1}{\sqrt{2}} \\ \frac{S_{12}}{\sqrt{2}} & -\frac{C_{12}}{\sqrt{2}} & \frac{1}{\sqrt{2}} \end{pmatrix}, \quad (26)$$

where $S_{12} \sim \frac{M_D}{M_N} = \frac{y_D v}{y_N v'}$ controls the mixing between ν_L and ν_R . Altogether:

$$\begin{pmatrix} \nu_l \\ \nu_h \\ \nu'_h \end{pmatrix} = U \begin{pmatrix} \nu_L \\ \nu_R \\ S_1 \end{pmatrix} \quad (27)$$

Neglecting intergenerational mixing, eq. (23) can be rewritten as

$$M_{\nu_l} \sim M_D M_N^{-1} M_{S_1} (M_N^T)^{-1} (M_D^*)^T = \sqrt{2} \frac{y_{s1} y_D^2}{y_N^2} \frac{v^2}{v'}. \quad (28)$$

To obtain light neutrino masses compatible with experiments (i.e., $M_{\nu_l} < 1$ eV), considering $v' = 10$ TeV, eq. (28) simplifies to

$$y_{s1} \left(\frac{y_D}{y_N} \right)^2 < 10^{-10}. \quad (29)$$

A possible solution, which allows for the heavy neutrinos in the $\mathcal{O}(100)$ GeV range (suitable for the LHC), requires $y_N \sim \mathcal{O}(10^{-2})$. Then, $y_D \sim \mathcal{O}(10^{-5})$ implies $y_{s1} < \mathcal{O}(10^{-3})$.

A fundamental difference with the model of Ref. [54] exists in the non-interacting neutrino sector, the S_2 fields (here we consider only one of them). In eq. (20) it is clear that the S_2 field acquires a mass after the $U(1)_{B-L}$ symmetry breaking when χ gets a vev (called v'):

$$M_{S_2} = \sqrt{2} y_{s2} v'. \quad (30)$$

Therefore, its mass is expected to be of $\mathcal{O}(1)$ TeV, if also v' is at TeV scale. The Z_2 symmetry protects this particle from decaying, making it a suitable dark matter candidate.

Given that the S_2 field acquires a mass through a Yukawa term similar to the Majorana mass of the S_1 field (see eq. (30)), it is expected that their Yukawas (y_{s1} and y_{s2}) lie in similar ranges. As we will see, the S_2 field is a suitable DM candidate if its mass is of the order of hundreds of GeV, that requires $y_{s2} \sim \mathcal{O}(10^{-2})$, similar to the order of magnitude for y_{s1} required from light neutrino mass constraints, as one would naively expect.

Such range in masses allows for a suitable relic density due to resonant decays via the heavy Higgs boson only. Relaxing the latter requirement ends up in a larger variety of phenomenological implications, such as heavier DM candidates and resonant annihilations with the Z' boson, as well as more interesting LHC phenomenology in the interacting neutrinos due to a larger mixing with the LH fields, because of a larger y^D coupling of $\mathcal{O}(10^{-2})$.

III. RESULTS: DARK MATTER

The S_2 field as defined in the previous section is a suitable candidate for WIMP-like cold dark matter because it is odd with respect to the Z_2 symmetry and thus stable on cosmological time scales. An abundance of S_2 particles which is thermally produced in the early universe will survive after freeze-out from the thermal bath until today to make up the observed dark matter in the universe.

To compute the relic density of S_2 dark matter we used the program MicrOMEGAs [61, 62], in which the model has been implemented via LanHEP [63]. The remaining numerical analysis has been performed in CalcHEP [64]. The a priori unknown Yukawa couplings of the extra neutrino fields turn out to be negligible in the calculation of the relic density. For concreteness, in the following we fixed $y_{\nu 3}^D = 10^{-4}$, $y_{\nu 3}^M = 0.06$ and $y_3^{S_1} = 10^{-5}$, owing to $m_{\nu l} = 1.5 \cdot 10^{-10}$ and $M_{\nu_h} = M_{\nu'_h} = 636$ GeV for $v' = 15$ TeV. Further, the Higgs mixing angle has been set to $\sin \alpha = 0.1$.

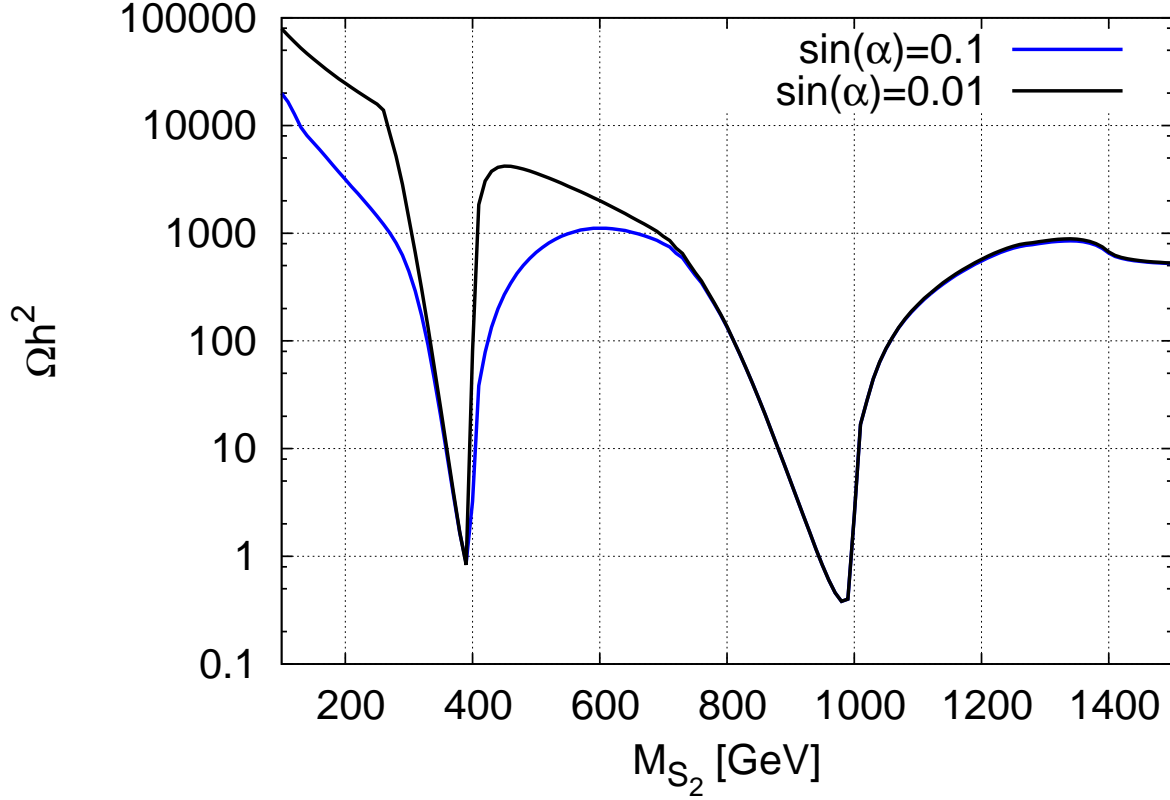


FIG. 2: DM relic abundance as a function of the candidate mass for $M_{H_2} = 800$ GeV and $M_{Z'} = 2$ TeV, for two different choices of the scalar mixing angle. For the other parameters, see the text.

The recent discovery at the LHC [65] if interpreted as a SM-like Higgs boson, does not constrain this parameter very strongly yet, but implies it is small. The neutrino couplings and scalar mixing angle α will stay fixed to these values throughout the following analysis. Therefor our free parameters are $M_{Z'}$, M_{H_2} , g_{BL} , g_2 and the dark matter mass M_{S_2} .

For a first assessment of the situation we fix $g_{BL} = 0.1$, $g_2 = 0$, $M_{H_2} = 0.8$ TeV and $M_{Z'} = 2.0$ TeV. We expect the coupling g_{BL} and the Higgs mixing angle α to have an impact on the relic density, as they control the interaction strength of the Z' boson and the Higgs bosons, respectively, with the DM particle and the remaining particles of the model.

We compute the relic density as a function of the S_2 mass, as shown in figure 2. Comparing

the resulting graph with the observed abundance of dark matter:

$$\Omega_{DM}h^2 = 0.1117^{+0.0053}_{-0.0055}, \quad (31)$$

it turns out that for $M_{S_2} \not\approx M_{Z'}/2$ and $M_{S_2} \not\approx M_{H_2}/2$, the relic density is always too large by several orders of magnitude. There are two ways to reduce Ωh^2 : one is to lower the vev v' (which results in an increase in the Yukawa couplings), which will make the heavy Higgs resonant annihilation more efficient. The other possibility is to increase the coupling g_{BL} to enhance the Z' annihilation efficiency. However, an arbitrarily small relic density is forbidden, due to the LEP lower limit on v' of eqs. (18)–(19), and by the upper limit on g_{BL} due to RGE analysis [66]. Obviously, the S_2 Yukawa coupling is linked to the v' and therefore to g_{BL} .

We conclude from this first assessment that for the S_2 particle to be the dark matter, it must annihilate efficiently via a resonant heavy Higgs or via the Z' boson. Also the precise value of α is not relevant in the resonant regions, that are the only important regions for our analysis, as long as it is small, as expected if one relies on the confirmation of the LHC hints on a light SM-like Higgs boson. In the following we will study these two different mechanisms, fixing $\sin \alpha = 0.1$.

A. Higgs boson resonant annihilation

For small DM masses, where the main annihilation channel is via the heavy Higgs boson, the relic density is proportional to the B - L -breaking vev, since $\Omega h^2 \propto \frac{1}{\sigma} \propto \frac{1}{y_{S_2}^2} \propto v'^2$, via eq. (30). According to LEP, the vev is constrained from below as in eqs. (18)–(19) for the models of interest. The left panel of figure 3 shows that for $M_{S_2} \simeq M_{H_2}/2$ the minimal abundance just matches the observed value.

In the right panel of figure 3 we see that the demand, that the resonant S_2 annihilation proceeds via the Higgs channel, strongly constrains from above the vev v' as a function of the heavy Higgs mass M_{H_2} , $v' \leq 11.7$ TeV. It is interesting to note that this upper bound on v' is independent of the scalar mixing angle α . At the same time, the heavy Higgs mass is constrained from above depending on the value of α . For instance for $\sin \alpha = 0.1$, $M_{H_2} \leq 1.8(3.0)$ TeV for $g_2 = 0(g_2^{min})$.

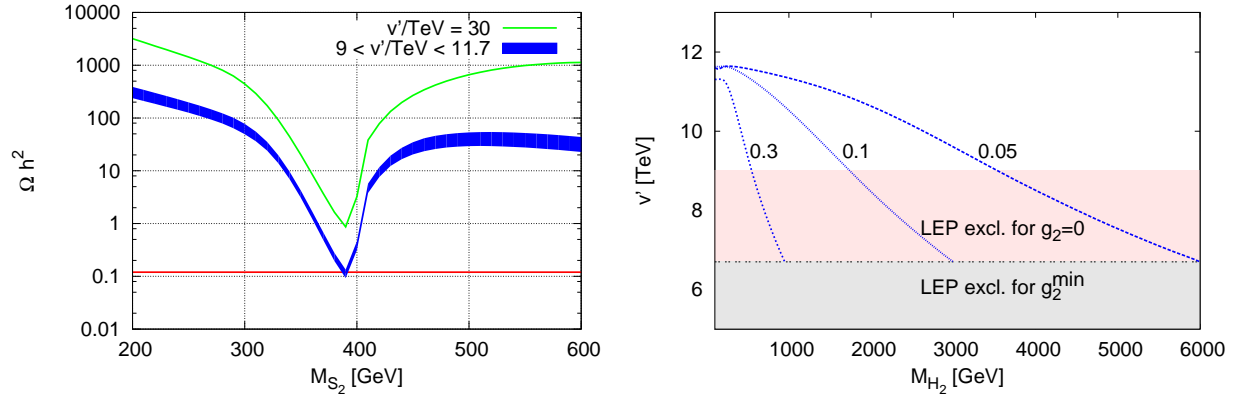


FIG. 3: (*Left*) DM relic abundance as a function of the DM mass for $M_{H_2} = 800$ GeV, for some values of the vev v' and $\sin \alpha = 0.1$. The blue shading represents values of the vev for which an allowed DM mass exists, for this M_{H_2} , as taken from the right panel.

(*Right*) Allowed region for v' for which a DM mass yields the correct relic density, as a function of the heavy Higgs boson mass, for the three values of $\sin \alpha = 0.05, 0.1, 0.3$. The LEP exclusion is as in eqs. (18)–(19).

B. Z' boson resonant annihilation

Due to LHC direct searches, the Z' boson mass has to be above $2 \div 2.5$ TeV, depending on the gauge mixing coupling g_2 . For the Z' annihilation mechanism to be effective, as already stated, a resonant condition has to be matched, meaning that the DM candidate has to be rather heavy. However, the resonance decay can still be very effective.

Figure 4 shows in more detail that only around the Z' resonance the annihilation is sufficient to match the abundance constraint. For simplicity, we fix here $g_2 = 0$, so that the relic density is inversely proportional to the square of the gauge coupling, i.e., $\Omega h^2 \propto \frac{1}{(g_{BL})^2}$. The lower limit on the vev v' translates into an upper limit on the coupling, for a fixed Z' mass: in figure 4 the Z' mass is 2 TeV, which gives the upper limit $g_{BL} < 1/3$, via eq. (17). Then, the demand $\Omega h^2 \simeq 0.1$ gives a lower limit for the gauge coupling, $g_{BL} > 0.21$.

Although the gauge coupling g_{BL} is expected to give the major contribution to the relic abundance evaluation, the impact of g_2 might be not negligible either. For fixed $M_{Z'}$ and g_{BL} values, we study here the effect of the mixing gauge coupling. We choose two different sets of $M_{Z'}$ and g_{BL} such that for $g_2 = 0$ the relic density just satisfies the abundance constraint. Figure 5 then shows that g_2 can have an impact on the relic abundance, especially it can

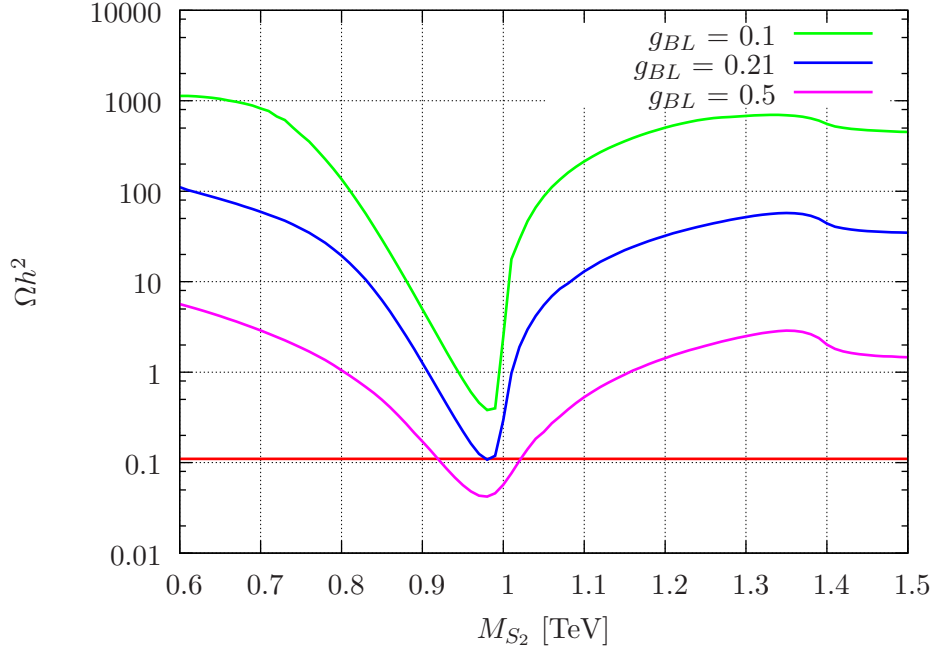


FIG. 4: Relic density as a function of the DM candidate mass around the Z' peak. The lower curve, for $g_{BL} = 0.5$, is for illustrative purposes only, being already excluded by LEP.

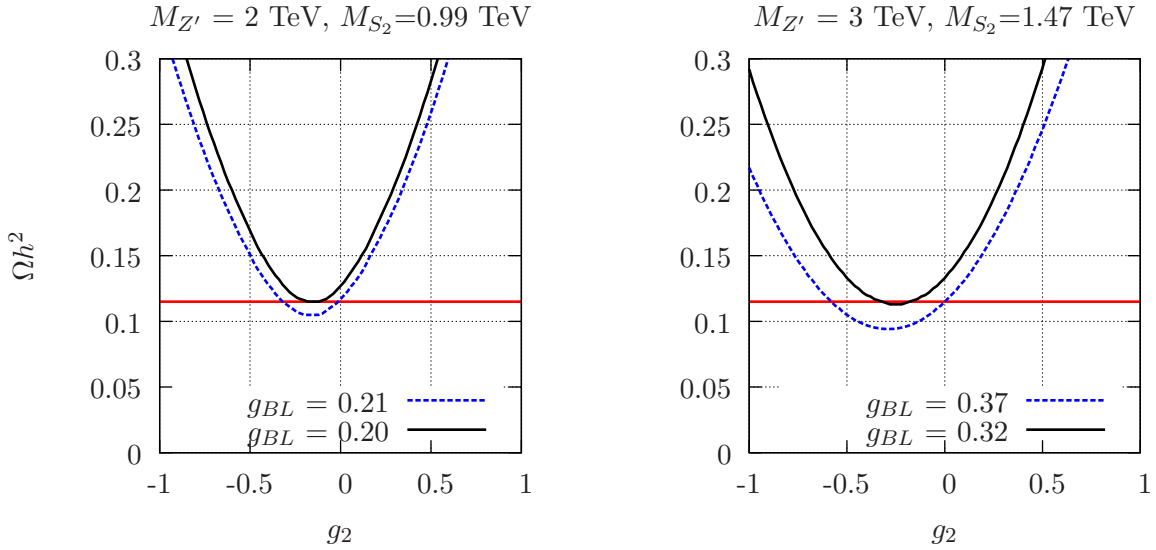


FIG. 5: Variation of the relic density (at the Z' resonance) with g_2 for a choice of g_{BL} (see the text for further details), for (*left*) $M_{Z'} = 2$ TeV and (*right*) $M_{Z'} = 3$ TeV.

lower the latter when assuming small negative values.

A general feature is that $\Omega(g_2)$ is growing for $|g_2| \rightarrow 1$, so that there exists a value of

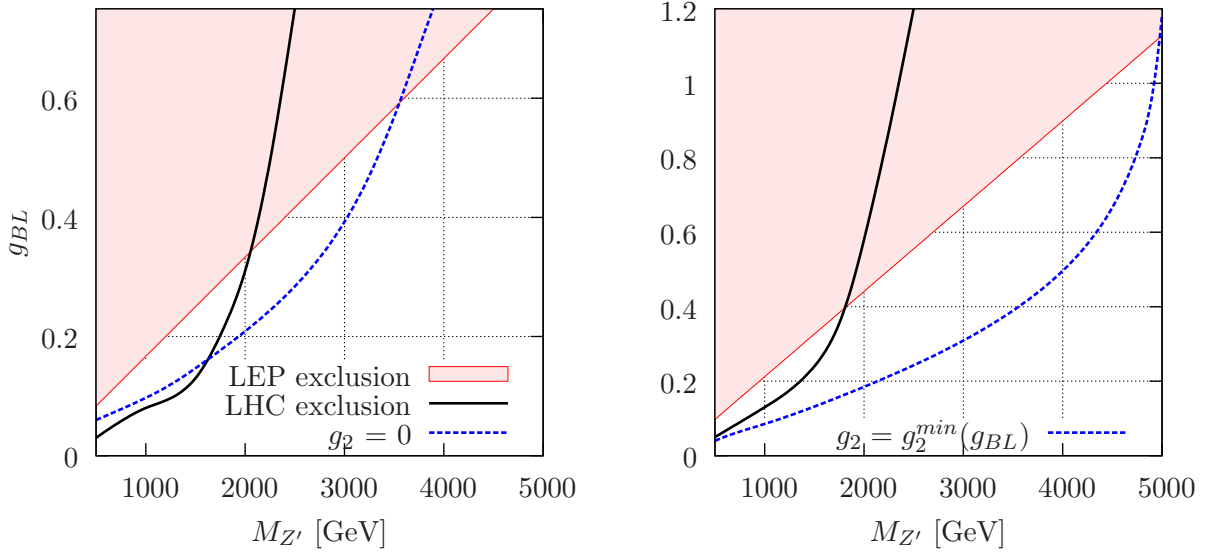


FIG. 6: Existence of a suitable DM candidate: the allowed region is the one above the dashed curves, in the $g_{BL} - M_{Z'}$ plane, (left) for $g_2 = 0$ and (right) for $g_2 = g_2^{\min}(g_{BL})$, which minimizes the Z' width, hence allowing for a smaller g_{BL} (and therefore a smaller cross section) per fixed $M_{Z'}$. The red shading combinations are forbidden by LEP (eqs. (18)–(19), respectively), the black (solid) lines are the LHC exclusion, as in table I.

g_2 that minimizes $\Omega(g_2)$. This value of g_2 corresponds to the g_2^{\min} of eq. (12), that also minimizes the Z' width. Fixing g_2 to $g_2^{\min}(g_{BL})$ returns a minimum value for g_{BL} such that the abundance constraint can be matched, which is roughly 5% lower than the lower limit on g_{BL} in the case $g_2 = 0$. We can now study the relic density as a function of the Z' boson mass. This is done in figure 6, that shows the range of allowed values for g_{BL} as a function of $M_{Z'}$. The curves, for both $g_2 = 0$ and $g_2 = g_2^{\min}(g_{BL})$, limit the existence of a DM candidate mass with suitable relic density above the curves themselves.

As a result, the Z' mass is constrained to $M_{Z'} \leq 3.5(5.0)$ TeV for the $g_2 = 0$ and the $g_2 = g_2^{\min}(g_{BL})$ case, respectively, which also imposes an upper bound on the gauge couplings when combined with the exclusion limits from LEP. Notice that this value for the Z' mass is within the LHC ultimate reach [41]. Since we must be on a resonant regime, the upper bound on the Z' mass translates on a upper bound for the S_2 mass, $M_{S_2} \leq 2.5$ TeV.

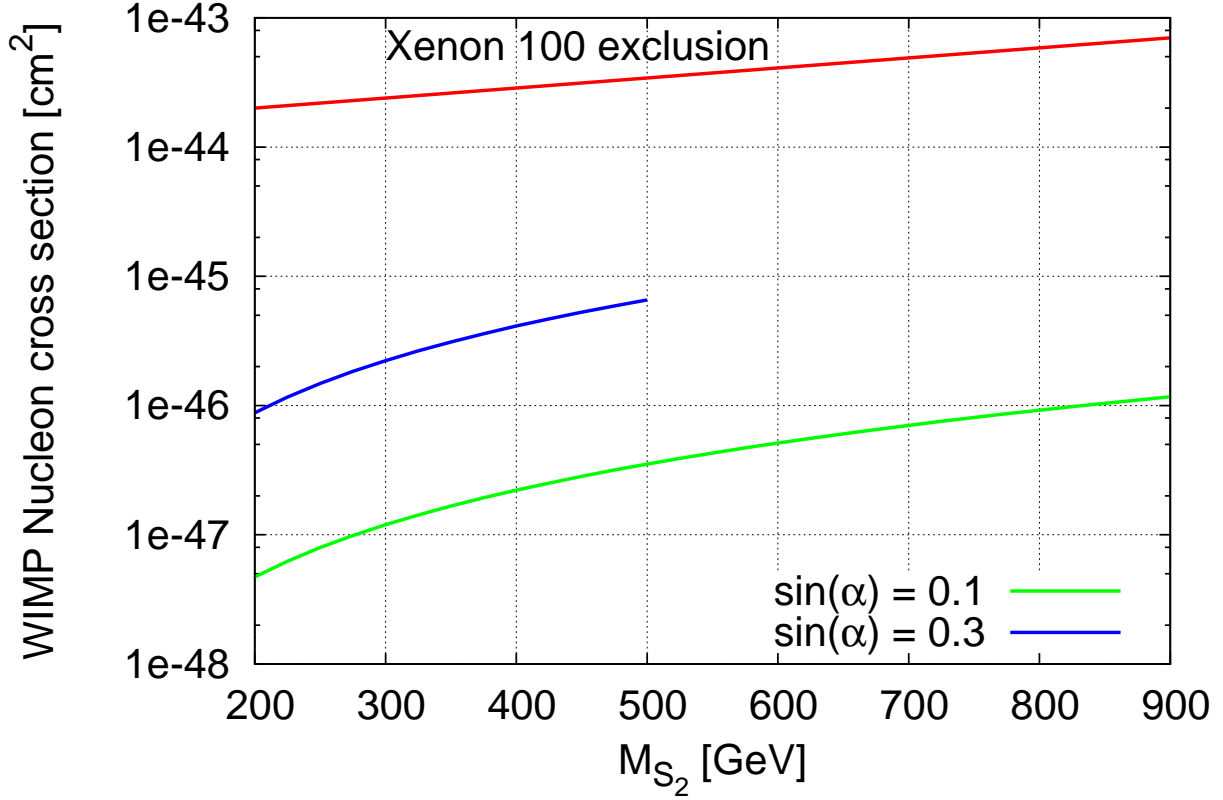


FIG. 7: Spin-independent direct searches for maximum Yukawa couplings $y_{S_2} \sim M_{S_2}/v'$. Only allowed masses are plotted, from figure 3(right), for $g_2 = g_2^{min}$.

C. Direct detection

The spin independent leptino-nucleon interaction is mediated by the light Higgs boson and is therefore sensitive to the value of $(\sin \alpha \cos \alpha)^2$. Figure 7 shows that even for the rather large value of $\sin \alpha = 0.3$, the leptino-nucleon cross section is several orders of magnitude below the actual exclusion limits from XENON100 [67], at most of roughly 10^{-44} cm² for a DM candidate mass of around 50 GeV. Since we know that $\sin \alpha$ has to be small, also future direct detection experiments will not be able to restrict or to detect the leptino. We checked that the exchange of Z' bosons consistent with LEP exclusion limits gives rise to smaller cross sections than Higgs exchange, hence we did not further consider this process. Even easier to evade are the constraints from spin-dependent leptino-nucleon interaction

experiments. Here the only mediator that can play a role is the Z' boson, too heavy for these cross sections to be of any interest.

D. Extension to N_ℓ families

As we will discuss in the next section, a successful leptogenesis needs $N_\ell \geq 2$ to provide the required large CP -violating phases. Here we want to comment on the impact of extra leptino generations in the DM analysis carried out so far. The extra neutrinos affect the Z' width only, see eq. (11), and will be discussed later. First we focus on the leptinos. If the S_2^i fields possess a mass hierarchy we have a situation similar to SUSY dark matter. The heavier particles decouple earlier and decay rapidly to the still thermalized lighter particles. An observable effect appears only if the relative mass difference is less than 10%. If they are mass degenerate, however, we get $\Omega_{N_\ell} = N_\ell \cdot \Omega_1$. Thus to have resonant annihilation the couplings have to be increased by $\sqrt{N_\ell}$. Figure 8 shows how this affects the parameter space. For resonant heavy Higgs annihilation the vev v' must be reduced by a factor $\sqrt{N_\ell}$, so that only the cases $N_\ell = 1, 2$ can still match the relic abundance for Higgs resonant annihilation, while $N_\ell = 3$ is just touching the LEP exclusion limits (for $g_2 = g_2^{min}$) from below and therefore such resonant mechanism is not sufficient.

The case of resonant Z' annihilation works well also for $N_\ell = 3$. However, the allowed parameter space becomes tighter for higher N_ℓ . These results are valid when all the leptinos are mass-degenerate, and represents the worst possible case. All others, i.e. when just 2 are degenerate or with a tight mass hierarchy, will be somewhere between the case of 1 generation and the case of 3 generations, degenerate in mass.

We turn now to study the effect of having 3 generations of leptinos, as required for leptogenesis, on the dark matter relic abundance due to the extra heavy neutrinos. The model with 3 generations of leptinos has 3 light and 6 heavy Majorana neutrinos in total. The one with only 1 generation of leptinos instead accounts for 3 light neutrinos, 2 of them being Dirac particles, and 2 heavy Majorana neutrinos. Although the total number of relativistic degrees of freedom, commonly addressed as g^* , is basically unchanged, the proliferation of neutrinos affects the Z' width, and this could impinge on the evaluation of the relic abundance at the Z' resonance. For the same setup as in section III, figure 9 shows the total Z' width in the 2 different cases.

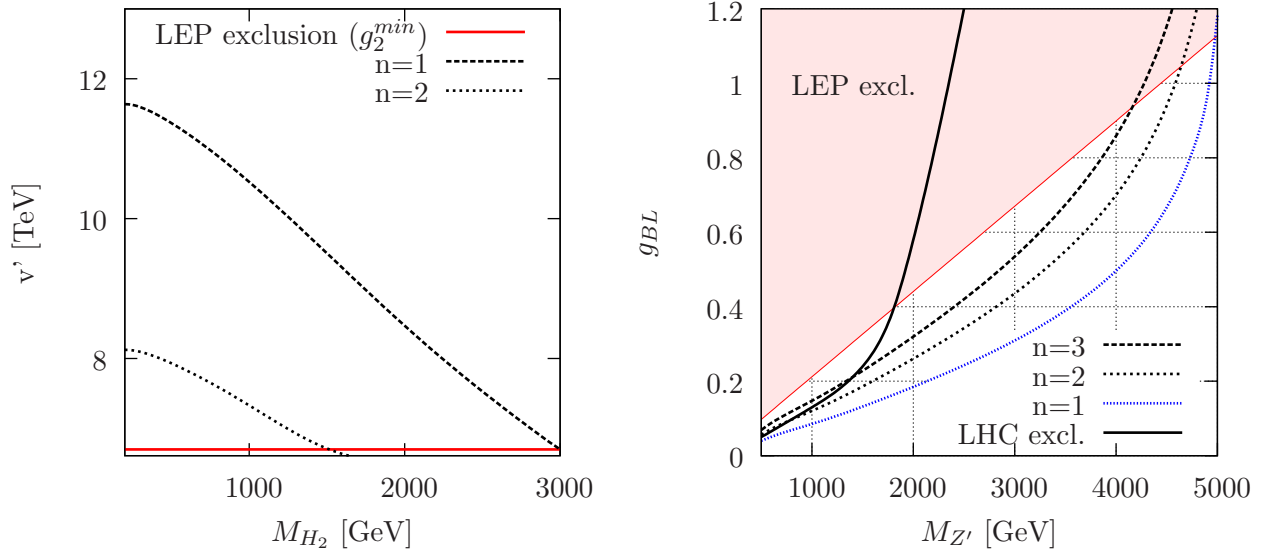


FIG. 8: (Left) $n \equiv N_\ell = 3$ is not shown as always disallowed. Here, $\sin \alpha = 0.1$. (Right) Allowed parameter range for resonant Z' annihilation for $n \equiv N_\ell = 2, 3$ families of leptinos.

The curves are for $g_2 = g_2^{\min}(g_{BL})$.

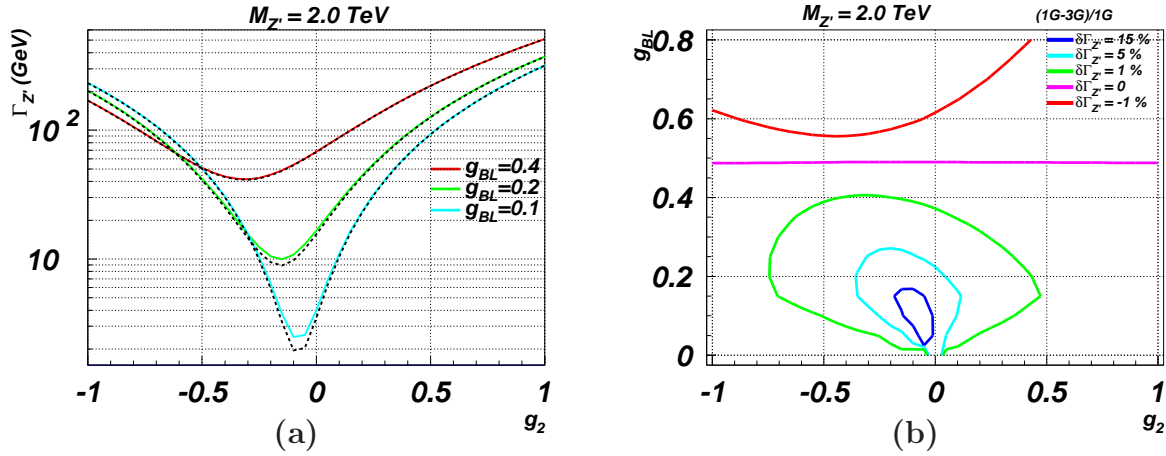


FIG. 9: (a) Total Z' width for selected g_{BL} values (solid lines refer to $N_\ell = 1$, dashed lines to $N_\ell = 3$) and (b) percentage variation between the 2 models in the $g_{BL} - g_2$ plane, for $M_{Z'} = 2.0$ TeV. For the heavy neutrino masses, see the text.

Despite the larger numbers of possible final states into which the Z' boson can decay for $N_\ell = 3$ with respect to $N_\ell = 1$, the latter case has a larger Z' boson partial width into light neutrinos since 2 of them are Dirac particles. Moreover, finite mass and threshold effects for the heavy neutrinos diminish their branching ratios. Altogether, this induces an almost

exact compensation between the larger number of heavy neutrinos for $N_\ell = 3$ and the larger partial width due to the 2 Dirac light neutrinos for $N_\ell = 1$. In the chosen setup for the Yukawa couplings, the relative variation of the total Z' width is below 1% in most of the parameter space, getting above 5% only in a limited region around small values of the gauge couplings. In these plots, heavy neutrino masses are $m_{h\nu} = 637$ GeV for $g_{BL} = 0.2$, and scale inversely with the latter. When their mass is above $M_{Z'}/2$ (or slightly before due to threshold effects, i.e., here for $g_{BL} < 0.15$), the $Z' \rightarrow NN$ ($N = \nu_h, \nu'_h$) channels are all suppressed or simply forbidden and the above mentioned compensation is not taking place. Even though the model with 1 generation of leptinos has here a bigger Z' width due to the larger partial widths into the light Dirac neutrinos, the excess is never above 20% of the total Z' width in the case of 3 generations, and only in a tiny corner of the parameter space, where both gauge couplings are small.

As intimated, the variation of the total Z' width is very small, mostly below 1%, so that the impact on the DM relic abundance is also negligible. This is because, as we can see by comparing figure 9(a) to figure 5(left) (by reading the value of the total width in the first figure and the respective value for the relic abundance in the second figure, for $g_{BL} = 0.2$ and $M_{Z'} = 2.0$ TeV, as a function of g_2), the relic abundance scales with the square root of the Z' width. A 1% variation in the total width determines a 0.5% variation of the relic density.

We can conclude that the DM study for $N_\ell = 1$ of section III is independent of the particular value of N_ℓ if the leptinos are not mass degenerate.

IV. FURTHER RESULTS

We comment here on further possible implications of the model. First, the possibility of leptogenesis is described, necessarily requiring the presence of more than one generation of leptinos. Next we comment on how our model has an effect on the effective number of light degrees of freedom N_{eff} , measured in cosmology. Extra degrees of freedom consist of the RH neutrinos, forming Dirac fermions with the LH counterparts when less than 3 generations of leptinos are present.

A. Leptogenesis

As discussed in the introduction, the inverse seesaw mechanism is a suitable mechanism for a large CP asymmetry even at the TeV scale due to the naturally small mass splitting between the heavier neutrino eigenstates, a result that is built in into the model and does not require fine tuning. Hence, it is a favourable model for the so-called resonant leptogenesis. Despite a large loop enhancement due to resonance, large phases are still required in the off-diagonal terms of y_D to obtain an $\mathcal{O}(1)$ CP asymmetry. In the $N_\ell = 1$ model discussed so far, the first 2 generations of light neutrinos require $\mathcal{O}(10^{-12})$ Yukawas, which in turn also means that the CP -violating decays of the heavy Majorana neutrino pair are similarly suppressed. Large phases but small masses can be achieved extending the seesaw mechanism to the other generations.

We will show that leptogenesis, compatible with neutrino masses and mixing, is possible in the model. The CP asymmetry is generated by the decays of the heavy neutrinos:

$$\varepsilon_i = \frac{1}{8\pi} \sum_{j \neq i} \frac{\text{Im} \left[(y_D y_D^\dagger)_{ij}^2 \right]}{\sum_\beta [y_{i\beta}^D]^2} f_{ij}^\nu, \quad (32)$$

requiring large phases in the off-diagonal elements of the Dirac Yukawa matrix y_D . In our model, these large phases, compatible with neutrino data, are possible for $N_l \geq 2$. The lepton number violating loop factor f_{ij}^ν , when quasi-degenerate heavy neutrino pairs are considered, is

$$f_{ij}^\nu = \frac{M_j^2 - M_i^2}{(M_j^2 - M_i^2)^2 + (M_j \Gamma_j - M_i \Gamma_i)^2}. \quad (33)$$

In the inverse seesaw case, $M_j \sim M_i$ and $\Gamma_j = \Gamma_i \equiv \Gamma$ are naturally recovered, so that f_{ij}^ν can easily be $\mathcal{O}(1)$. Notice that no fine tuning is required.

Once an asymmetry in the lepton sector is produced, electroweak sphaleron processes take place and move the asymmetry to the baryon sector

$$\eta_B = \frac{28}{79} \eta_{B-L}. \quad (34)$$

Altogether, the final baryon asymmetry can be written as

$$\eta_B \sim 10^{-2} \varepsilon_i \kappa_i (z \rightarrow \infty), \quad (35)$$

where the 10^{-2} pre-factor accounts for the sphaleron efficiency and for photon dilution after recombination. The CP asymmetry ε_i was defined in eq. (32), while κ_i is the efficiency

factor obtained after solving the relevant Boltzmann equations:

$$\kappa_i(z) \sim \int_{z_0}^z dz' \frac{dN_{N_i}^{eq}(z')}{dz'} \frac{D(K_i, z')}{D(K_i, z') + 4S_{Z'} N_{N_i}^{eq}(z')} \times \exp \left[- \int_{z'}^z dz'' W_{ID}(K_i, z'') \delta_i^2 \right], \quad (36)$$

with $\delta_i = \frac{|M_i - M_j|}{\Gamma} \ll 1$ suppresses considerably the inverse decay (ID) wash-out.

All the quantities are defined in Ref. [53]. Particularly important is the $S_{Z'}$ term, the Z' scattering processes, which induces a wash-out of the final asymmetry. It has been verified that these Z' -induced wash-out processes do not have a severe impact on the final baryon asymmetry. This is due to the nature of the inverse seesaw mechanism, that allows for large Yukawa couplings, overcoming the Z' processes. A similar conclusion was reached in Ref. [46], in which the authors explicitly showed that Z' processes do not spoil the leptogenesis if the Dirac Yukawa couplings are sufficiently large, as in the inverse seesaw case under examination.

In conclusion, following the similar case studied in Ref. [53], the Dirac Yukawa matrix y_D contains all information required to study the leptogenesis, entering both in ε_i and in Γ . It is sufficient that some off-diagonal elements in y_D are large and complex for leptogenesis to be possible. For definiteness, we have verified that the choice of the matrix in Ref. [53], that for $N_l = 3$ is a possibility in our setup, does yield the correct baryon asymmetry in our model. We stress again that a similar choice for y_D compatible with a successful leptogenesis is possible also for $N_l = 2$, even though only one off-diagonal element will be large in this case. This is in fact identical to $N_l = 3$ in the 1-flavour approximation.

Hence, we have proven that a successful leptogenesis, compatible with the observed pattern of neutrino masses and mixing angles is possible in our model when at least two generations of leptinos are present. However, the complete analysis of the leptogenesis in our model, as for instance the detailed comparison of $N_l = 2$ and $N_l = 3$ cases, or the impact of flavour effects outside the simple 1-flavour analysis, is beyond the scope of this paper and is left for future work.

B. Impact on N_{eff}^ν

As a last application of our model, we describe the possible implications on N_{eff}^ν . This observable counts the relativistic energy content in the universe at the time of the last scattering surface in terms of an effective number of neutrino species. Indications from

cosmology result in the observed value of $N_{eff}^{CMB} = 4.56 \pm 0.75$ when combining WMAP, the Atacama Cosmology Telescope, baryonic acoustic oscillations data and the measurement of the Hubble parameter H_0 [68]. Notice that the SM LH neutrinos only would yield $N_{eff}^\nu \sim 3$. This mismatch is typically interpreted as an indication of the existence of some extra relativistic degrees of freedom that effectively contribute as one unit of N_{eff}^ν . In our setup, the only extra degrees of freedom that can be relativistic at the last scattering surface are the RH neutrinos, when they have only Dirac mass terms, i.e., for less than 3 generations of leptinos, as described in section IIID. On the other hand, as we have seen previously, a successful leptogenesis requires more than one generation of leptinos. If one would like to explain both leptogenesis and ΔN_{eff}^ν in our model, less than three generations of leptinos should be considered, given that $N_\ell = 3$ leads to an effective number $N_{eff}^\nu \sim 3$ as in the standard model.

An estimate of the impact on N_{eff}^ν is done as follows. First, the decoupling temperature and the contribution to N_{eff}^ν depend on the parameters $m_{Z'}$, g_{BL} , g_2 , given that only the Z' boson can keep the RH neutrinos in thermal equilibrium with electrons, as the RH neutrinos have vanishing hypercharge. Naively, the bigger the Z' cross section, the lower the decoupling temperature, which in turns also means that less degrees of freedom are relativistic at decoupling. Overall, this increases the RH neutrino contribution to N_{eff}^ν . Close to LEP exclusion limits, we get a minimum decoupling temperature of 380 MeV. At this temperature the number of relativistic degrees of freedom is 20 if the QCD phase transition takes place at 450 MeV [69]. This yields the highest value for ΔN_{eff}^ν in our model: $0.18(3 - N_\ell)$. For lower QCD phase transition temperatures, the numbers of relativistic degrees of freedom rapidly increases, suppressing the impact on N_{eff}^ν to the percent level or below.

Thus we see that although the model can contribute to N_{eff}^ν , its impact is only marginal, especially if one would want to implement a successful leptogenesis. In fact, $N_\ell \geq 2$ is required by latter, while $N_\ell \rightarrow 0$ maximises ΔN_{eff}^ν .

V. CONCLUSIONS

We constructed a simple model in the class of the minimal Z' models, where the extra $U(1)$ gauge field is coupled only to hypercharge and B-L. We added right-handed neutri-

nos for each generation of fermions, as required by the absence of chiral anomalies in the theory. Beyond this we added extra pairs of leptons with fractional lepton number, which we therefore called leptinos. One of the leptino in the pairs is chosen to be even and the other one to be odd under an additional Z_2 charge. We were able to construct an inverse seesaw mechanism for neutrino masses. The mechanism is natural in the sense that all possible terms consistent with the symmetries of the theory are present in the Lagrangian. Nonetheless some of the entries in the neutrino mass matrix are zero, because of the choice of the representations. The reason is the presence of a fractional lepton charge, which is the new feature of the model.

The odd leptino is a candidate for dark matter, as it is weakly interacting, massive and stable. We have shown that the correct dark matter density can be generated if the leptino is annihilated through a resonance by either the Z' -boson or the Higgs-boson, related to the breaking of the B-L symmetry. We studied limits on the parameters of the theory, coming from LEP, hadron colliders and the dark matter abundance. We found that the limits are such, that the Z' -boson and extra Higgs-boson lie within the range of the LHC, though the full design luminosity might be needed. The cross sections are too small for present direct search experiments for dark matter. Resonant leptogenesis is possible in the presence of more than one pair of leptinos, but contains no particularly new features compared to other models of leptogenesis.

In conclusion the model provides a very simple extension of the standard model, containing a number of desirable features like dark matter, leptogenesis and (inverse) seesaw. At the same time, being a singlet extension, it does not lead to phenomenological problems, such as flavour changing neutral currents.

The model we constructed appears to be able to give a realistic description of cosmological data and can be tested at the LHC. However, it is not possible to describe at the same time leptogenesis and N_{eff}^ν .

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